

Nonuniform target illumination in the deflagration regime: Thermal smoothing

J. Sanz, J. A. Nicolás, J. R. Sanmartín, and J. Hilario

Escuela Técnica Superior de Ingenieros Aeronáuticos, Universidad Politécnica, 28040 Madrid, Spain

(Received 30 July 1987; accepted 29 February 1988)

Thermal smoothing in the plasma ablated from a laser target under weakly nonuniform irradiation is analyzed, assuming absorption at n_c and a deflagration regime (conduction restricted to a thin quasisteady layer next to the target). Magnetic generation effects are included and found to be weak. Differences from results available in the literature are explained; the importance of the character of the underdense flow at uniform irradiation is emphasized.

I. INTRODUCTION

Direct-drive laser fusion has usually been considered too sensitive to nonuniformities, e.g., in laser-energy deposition, and to instabilities at or outside the ablation surface. An alternative approach, one based on Hohlraum x rays, has been considered to bypass that difficulty. Recently, however, apparently moderate Rayleigh–Taylor growth rates and improved optics such as the induced spatial incoherence technique have made the direct drive more attractive.¹ It is thus essential to ascertain the physics and scaling involved: first, in any instability present in the corona blowing off the target, and, second, in the smoothing of nonuniformities in the driving. The latter aspect is the subject of the present paper.

In the direct-drive approach, the pressure needed to compress the target is generated by heat conducted from the critical to the ablation surfaces. It was recognized early on^{2,3} that conduction itself could result in smoothing, an effect (cloudy-day effect) first analyzed by Gardner and Bodner.⁴ The motion of the plasma was later included in an analysis by Manheimer *et al.*⁵ Evans⁶ carried out some time-dependent calculations and Bell and Eperlein⁷ took into account magnetic effects. Here we consider thermal smoothing within a well-specified model, taking special care about the location of the critical surface and the plasma flow at and below the critical density n_c . We obtain complete results and discuss how they are dependent on the character of that flow. Effects caused by underdense refraction of light, which are also important,⁸ will be considered in a future paper.

We assume the following: an almost uniform irradiation, absorption occurring only at n_c , and a deflagration regime (heat conduction restricted to a quasisteady layer enclosing the critical surface). (In Sec. V we will discuss how these same simplifications were used in Ref. 7 and their effect on the results.) If refraction is ignored, a study of the layer will suffice to determine the smoothing. Such a layer exists if the distance x_c between the critical and ablation surfaces is small compared with both the overall length of plasma L and its transverse dimension R (radius of a pellet or radius of the focal spot on a foil). In this regime $x_c \sim \bar{K} I_0^{4/3} \bar{m}^{7/2} / \rho_c^{7/3}$ and $L \sim R$ or $L \sim (I_0 / \rho_c)^{1/3} \tau$, whichever is smaller; $\bar{K}(Z_i)$ is the factor (weakly dependent on density and temperature through a Coulomb logarithm) in Spitzer's conductivity, Z_i and $Z_i \bar{m}$ are the ion charge number and mass, and $\rho_c \equiv \bar{m} n_c$. For a peaked pulse, τ and I_0 are the duration and peak-ab-

sorbed intensity; in many numerical calculations a pulse of constant intensity I_0 is used and τ is then the time at the end of the run. Note that in the large isentropic region outside the layer, unsteady and/or three-dimensional effects must be important.

In Sec. II we present the equations of the model. Sections III and IV develop the mathematics of the solution to zero and first order in the nonuniformity. A discussion of the results is given in Sec. V.

II. BASIC EQUATIONS

The equations for the conducting layer are those of a quasisteady, quasineutral, two-fluid plasma:

ion continuity,

$$\nabla \cdot \mathbf{v} n / Z_i = 0, \quad (1)$$

electron and total momentum,

$$0 = -\nabla(nT) - en(\mathbf{E} + \mathbf{v} \wedge \mathbf{B}/c) + \mathbf{R}, \quad (2)$$

$$\bar{m} n \mathbf{v} \cdot \nabla \mathbf{v} = -\nabla(nT), \quad (3)$$

and electron internal energy (or entropy),

$$\frac{3}{2} n (\mathbf{v} + \mathbf{u}) \cdot \nabla T - T (\mathbf{v} + \mathbf{u}) \cdot \nabla n = -\nabla \cdot \mathbf{q} - \mathbf{u} \cdot \mathbf{R} + I', \quad (4)$$

together with Ampère and Faraday laws,

$$\nabla \wedge \mathbf{B} = -(4\pi/c) e n \mathbf{u} \quad (\nabla \cdot \mathbf{B} = 0), \quad (5)$$

$$\nabla \wedge \mathbf{E} = 0. \quad (6)$$

Here n and T are the electron density and temperature, n/Z_i is ion density, and \mathbf{v} and $\mathbf{v} + \mathbf{u}$ are ion and electron flow velocities. Electron continuity follows from (1) and (5). We have neglected electron inertia and viscous terms, and considered a large Z_i . This justifies our neglecting ion pressure in (3), and direct heating (due to temperature relaxation) in (4), thus uncoupling system (1)–(6) from the ion energy equation. Here I' is a δ function term representing laser heating at n_c .

The ion friction on the electrons and the electron heat flux are given by⁹

$$\begin{aligned} \mathbf{R} = & -\alpha_0 (m_e n / \tau) \mathbf{u} - \beta_0 n \nabla T \\ & - (\beta_0'' / \delta_0) n (\tau e \mathbf{B} / m_e c) \wedge \nabla T, \end{aligned} \quad (7)$$

$$\begin{aligned} \mathbf{q} = & -\bar{K} T^{5/2} \nabla T + \beta_0 T n \mathbf{u} \\ & - (\gamma_0'' / \gamma_0 \delta_0) \bar{K} T^{5/2} (\tau e \mathbf{B} / m_e c) \wedge \nabla T. \end{aligned} \quad (8)$$

The first terms above represent simple Ohm's and Fourier laws, while those at the end represent Nernst and Righi-Leduc effects, respectively. The terms in the middle are the thermoelectric effects. Also τ is an electron collision time, $\tau \equiv m_e \bar{K} T^{3/2} / \gamma_0 n$. The factors in (7) and (8) are functions of Z_i . For Z_i large, $\alpha_0 \approx 0.295$, $\beta_0 \approx 0.877$, $\gamma_0 \approx 12.5$, $\gamma_0'' \approx 10.2$, $\delta_0 \approx 0.0961$.⁹

For a one-dimensional layer, \mathbf{u} and \mathbf{B} vanish, \mathbf{E} is irrotational and given by (2), and Eqs. (1), (3), and (4) suffice to determine the n , \mathbf{v} , and T variables. For a weakly two-dimensional layer, as assumed below, \mathbf{u} and \mathbf{B} would be small, first-order quantities related by Eq. (5), and generated by density and temperature cross gradients through the requirement $\nabla \wedge \mathbf{E} = 0$ used in (2). Hall and Ettinghausen effects have not been included in (7) and (8), respectively, because they are proportional to the product $\mathbf{u} \wedge \mathbf{B}$, and thus are second-order quantities. For the same reason, a term $-enu \wedge \mathbf{B}/c$ was omitted from both momentum equations above; the \mathbf{v} terms on the left-hand side of (4) may then be written as $\nabla \cdot n\mathbf{v} (\frac{1}{2}\bar{m}\mathbf{v}^2 + \frac{5}{2}T)$. The factors in all terms of Eqs. (7) and (8) are, in general, even functions of $|\mathbf{B}|$, so to the order considered we used their values for vanishing $|\mathbf{B}|$.

Both the magnetic field and current affect the electron energy balance. The thermoelectric effects drop off the heating rate $-(\nabla \cdot \mathbf{q} + \mathbf{u} \cdot \mathbf{R})$, so that only Fourier (Spitzer) and Righi-Leduc terms, dependent on the temperature gradient, enter that rate. On the other hand, the electron entropy is convected with the velocity $\mathbf{v} + \mathbf{u}$, and thus there is a direct current effect on the electron heating (in Ref. 7, the term $-Tu \cdot \nabla n$ was missing from the energy equation). In the equation resulting from (2) and (6), \mathbf{u} enters through Ohm's term and \mathbf{B} through the $\mathbf{v} \wedge \mathbf{B}$ and Nernst terms; the thermoelectric effect drops off again (the term $-\beta_0 \nabla T$ in \mathbf{R}/n is irrotational).

III. UNIFORM IRRADIATION

We first review the results for a one-dimensional layer. Light impinges from the right on a dense, heavy target. Because of this, we may take the base of the corona (the ablation surface) to lie throughout at $x = 0$, and let $n \rightarrow \infty$ as $x \rightarrow 0$.¹⁰ The intensity absorbed at the critical surface is assumed to be transversely uniform, $I_c = I_{c0}$. Then, $I' = I_{c0} \delta(x - x_{c0})$ with $n(x_{c0}) = n_c$, and the solution depends only on x . We have $\mathbf{u} = \mathbf{B} = 0$, $\mathbf{q} = -\bar{K} T^{5/2} \nabla T$, and $\mathbf{R} = -\beta_0 n \nabla T$; \mathbf{E} is given by $e\mathbf{E} = -\beta_0 \nabla T - n^{-1} \nabla(nT)$ and Eqs. (1), (3), and (4) become

$$\frac{d}{dx}(n_0 v_0) = 0, \quad \frac{d}{dx}(\bar{m} n_0 v_0^2 + n_0 T_0) = 0, \quad (9)$$

$$\frac{d}{dx} \left[n_0 v_0 \left(\frac{1}{2} \bar{m} v_0^2 + \frac{5}{2} T_0 \right) - \bar{K} T_0^{5/2} \frac{dT_0}{dx} \right] = I' \equiv I_{c0} \delta(x - x_{c0}). \quad (10)$$

The above system determines the structure of the conducting layer, where x/x_{c0} is of order unity. As seen in the Introduction, the system and its solution fail at large enough x/x_{c0} , where either three-dimensional or unsteady effects, or both, are important and where, on the other hand, gradients are so weak that conduction is negligible and the flow isentropic. This is a typical problem with two widely different length scales, x_{c0} and L . An asymptotic expansion in the small parameter x_{c0}/L will be singular, and a method such as inner and outer expansion must be used to solve the entire problem. At large x/x_{c0} the layer's (inner) solution must match the isentropic (outer) solution at small x/L .¹¹

For a variety of situations the layer's exhaust velocity, v_0 (x/x_{c0} large), was proved equal to the local isentropic sound speed (Chapman-Jouguet condition).^{12,13} The proof was based on two facts: (i) system (9) and (10) does not

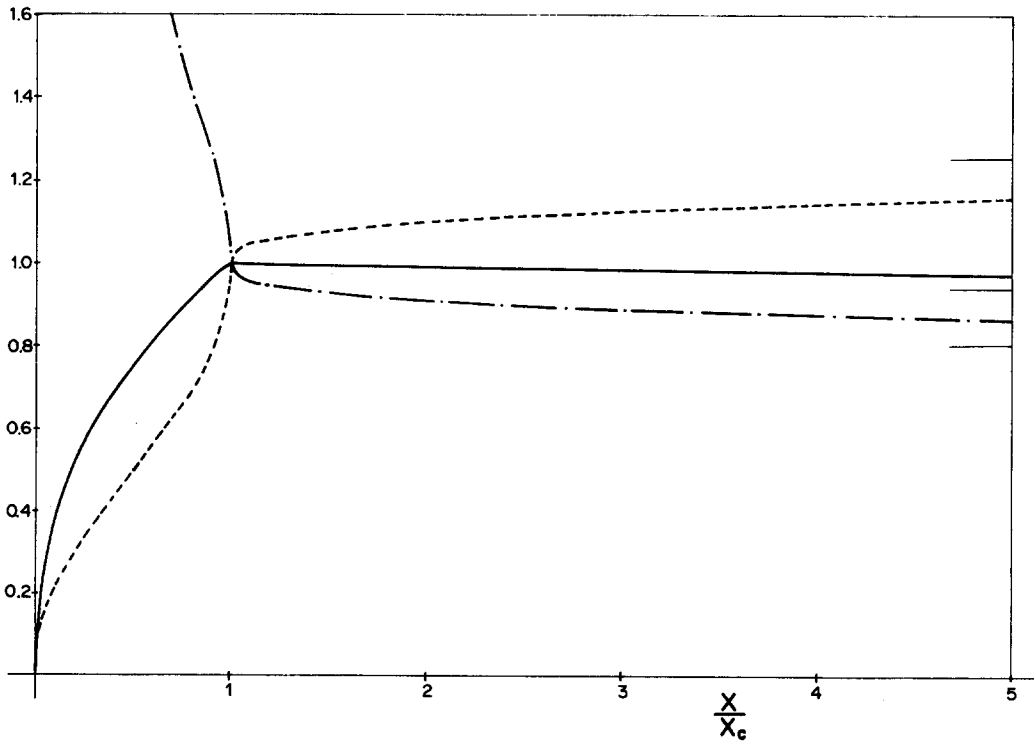


FIG. 1. Normalized fluid profiles in a deflagration layer under uniform irradiation; temperature —, density — —, velocity — · —. Absorption around the critical density at x_c ; target at $x = 0$. Within the scale of the layer, the profiles become constant as $x/x_c \rightarrow \infty$.

allow an isentropically supersonic exhaust velocity; and (ii) the flow outside, if it is quasisteady and spherical [$L \sim R < (I_0/\rho_c)^{1/3}\tau$]¹² or planar and unsteady though self-similar [$L \sim (I_0/\rho_c)^{1/3}\tau < R$],¹³ cannot start at (isentropically) subsonic speeds. (In combustion theory, the determination of the structure of a thin detonation or deflagration layer requires similar information on the flow outside.) The δ function at x_{c0} then requires that v_0 and the local isothermal sound speed be equal at the critical density; both n_0 and v_0 present infinite gradients there.

We then find that $n(x/x_{c0} \rightarrow \infty) \rightarrow 4n_c/5$; the entire structure of the layer can be determined and is shown in Fig. 1. Some quantities of interest such as the eigenvalue x_{c0} , the pressure P_{a0} at the ablation surface $x = 0$, the mass ablation rate $\bar{m}n_0v_0 = \text{const} = \dot{m}_{a0}$, or the temperature T_{c0} at the critical surface are^{12,13}

$$\begin{aligned} x_{c0} &\simeq 0.0326\bar{K}\bar{m}^{7/2}I_{c0}^{4/3}/\rho_c^{7/3}, \\ P_{a0} &= 8 \times 5^{-4/3}\rho_c^{1/3}I_{c0}^{2/3}, \quad \dot{m}_{a0} = \frac{25}{32}(P_{a0}^2/I_{c0}), \\ \frac{1}{5}T(x/x_{c0} \rightarrow \infty) &= T_{c0} = \bar{m}v_{c0}^2 = P_{a0}/2n_c. \end{aligned}$$

The more complex case of arbitrary Z_i has been solved in the past, also; for instance, if $Z_i = 1$, we have $n \rightarrow 0.61n_c$ as $x/x_{c0} \rightarrow \infty$ and $1.27T(x/x_{c0} \rightarrow \infty) = T_{c0}$.¹³

IV. WEAKLY NONUNIFORM IRRADIATION

We now consider a weakly nonuniform absorbed intensity

$$I_c = I_{c0} + I_{c1} \exp(iky),$$

with I_{c1}/I_{c0} small and k arbitrary. The laser heating term I' is then $I_c \delta(x - x_c)$, where x_c must be expanded in the form

$$x_c = x_{c0} + x_{c1} \exp(iky) + \dots \quad (11)$$

A naive expansion of the dependent variables would produce derivatives of the δ function for first and higher orders in the expansion. Worse than this, the sonic character of the zeroth-order solution at x_{c0} would result, for instance, in the perturbed density being infinite there. This would also make the numerical treatment around x_{c0} troublesome. A simple way to deal with this difficulty is to strain the x coordinate in the form

$$x = s[1 + (x_{c1}/x_{c0})e^{iky} + \dots], \quad (12)$$

that is, to define new variables $s \equiv x[1 - \exp(iky)x_{c1}/x_{c0} + \dots]$ and $y \equiv y$. Equations (11) and (12) show that $s = x_{c0}$ at $x = x_c$ to all orders; the heating term is then $I' = [I_{c0} + (I_{c1} - I_{c0}x_{c1}/x_{c0})\exp(iky) + \dots]\delta(s - x_{c0})$.

Variables as well as partial derivatives are now expanded to first order: $n = n_0(s) + n_1(s)\exp(iky)$ (similarly for T_0 and v_x), $v_y = v_{y1}(s)\exp(iky)$ (similarly for u_x , u_y , and B_z), and

$$\frac{\partial}{\partial x} = \left(1 - \frac{x_{c1}}{x_{c0}}e^{iky}\right)\frac{\partial}{\partial s}, \quad \frac{\partial}{\partial y} = \frac{\partial}{\partial y} - \frac{x_{c1}}{x_{c0}}ikse^{iky}\frac{\partial}{\partial s}.$$

The zeroth-order solution $n_0(s)$, $v_{x0}(s)$, and $T_0(s)$ is obviously given by the same functions of x (n_0 , v_0 , and T_0) determined by Eqs. (9) and (10).

To next order Eqs. (1), (3), and (4) yield a fifth-order linear system of ordinary differential equations with variable

coefficients, for n_1 , v_{x1} , v_{y1} , and T_1 :

$$\frac{d}{ds}(v_0n_1 + n_0v_{x1}) + ikn_0v_{y1} = 0, \quad (13a)$$

$$\frac{d}{ds}\left(2\dot{m}_{a0}v_{x1} + P_{a0}\frac{n_1}{n_0} + n_0T_1\right) + ik\dot{m}_{a0}v_{y1} = 0, \quad (13b)$$

$$\dot{m}_{a0}\left(\frac{dv_{y1}}{ds} + ik\frac{x_{c1}}{x_{c0}}s\frac{dv_0}{ds}\right) + ik(n_0T_1 + T_0n_1) = 0, \quad (13c)$$

$$\begin{aligned} (n_0v_{x1} + v_0n_1)\frac{d}{ds}\left(\frac{1}{2}\bar{m}v_0^2 + \frac{5}{2}T_0\right) \\ + \frac{d}{ds}v_0\left(\frac{5}{2}n_0T_1 + \dot{m}_{a0}v_{x1}\right) + \left(k^2 - \frac{d^2}{ds^2}\right)\bar{K}T_0^{5/2}T_1 \\ - I_{c1}\delta(s - x_{c0}) + \frac{x_{c1}}{x_{c0}}\left(\frac{d}{ds} - k^2s\right)\bar{K}T_0^{5/2}\frac{dT_0}{ds} \\ = \frac{\tau_0 dT_0/ds}{m_e v_0}\left(\frac{3}{2} - \frac{T_0 dn_0/ds}{n_0 dT_0/ds} + \Lambda_0^2\right)Q. \end{aligned} \quad (13d)$$

Here $Q \equiv \dot{m}_{a0}m_e u_{x1}/\bar{m}\tau_0 \propto B_z/n_0\tau_0$ is a heating rate, its equation being obtained from (2), (5), and (6):

$$\begin{aligned} \left(\frac{d^2}{ds^2} - k^2\right)Q + \frac{d}{ds}\left\{\frac{Q}{T_0}\frac{dT_0}{ds}\left[\frac{3}{2} \right. \right. \\ \left. \left. + \frac{\Lambda_0^2}{\alpha_0\gamma_0''}\left(\beta_0'' - \frac{\delta_0 m_e v_0}{\tau_0 dT_0/ds}\right)\right]\right\} \\ = -k^2\frac{v_0}{\alpha_0}\left(n_1\frac{dT_0}{ds} - T_1\frac{dn_0}{ds}\right). \end{aligned} \quad (14)$$

On the right-hand side of (13d) the Λ_0^2 term is the Righi-Leduc effect, the rest coming from the u terms on the left of (4). The Λ_0^2 terms in (14) are the $\mathbf{v} \wedge \mathbf{B}$ and Nernst effects, the rest on the left coming from Ohm's law. The dimensionless ratio $\tau_0(dT_0/ds)/m_e v_0$, where $\tau_0 \equiv \tau(n_0, T_0)$, grows monotonically from 0.20 at $s = 0$ to 0.24 at $s = x_{c0}^-$; it increases monotonically from -0.01 at $s = x_{c0}^+$ to zero as $s/x_{c0} \rightarrow \infty$. We also have

$$\Lambda_0^2 \equiv \frac{\gamma_0''}{\delta_0\gamma_0}\frac{\bar{K}T_0^{5/2}\tau_0}{n_c}\left(\frac{2\pi}{\lambda}\right)^2 \equiv \Lambda_c^2\left(\frac{T_0}{T_{c0}}\right)^4\frac{n_c}{n_0}, \quad (15a)$$

$$\begin{aligned} \Lambda_0^2 \simeq 0.29\left(\frac{\bar{m}}{2m_p}\right)^{4/3}\left(\frac{I_0}{10^{13}\text{ W cm}^{-2}}\right)^{8/3} \\ \times \left(\frac{\lambda}{0.26\text{ }\mu\text{m}}\right)^{22/3}\left(\frac{100}{Z_i \ln \Lambda}\right)^2 \end{aligned} \quad (15b)$$

(m_p = proton mass, $\ln \Lambda$ = Coulomb logarithm, λ = laser wavelength).

Eight conditions are clearly required to determine the solution to this seventh-order system with one eigenvalue x_{c1} . They are as follows.

(1) Zero current off the target (which otherwise would become charged), $QT_0^{3/2} \propto n_0u_{x1} = 0$ at $s = 0$.

(2) and (3) Pressure and mass flow rate bounded at the ablation surface, leading to $v_{x1} = T_1 = 0$ at $s = 0$.

(4) v_{y1} continuous at the ablation surface and hence $v_{y1} = 0$ at $s = 0$.

(5) $n_1 = 0$ at $s = x_{c0}$ (note that $n - n_0 = n_1e^{iky} + \dots$ and that at $s = x_{c0}$ both n , to all orders, and n_0 are equal to n_c).

(6) Fluid variables bounded at $s = x_{c0}$. This condition

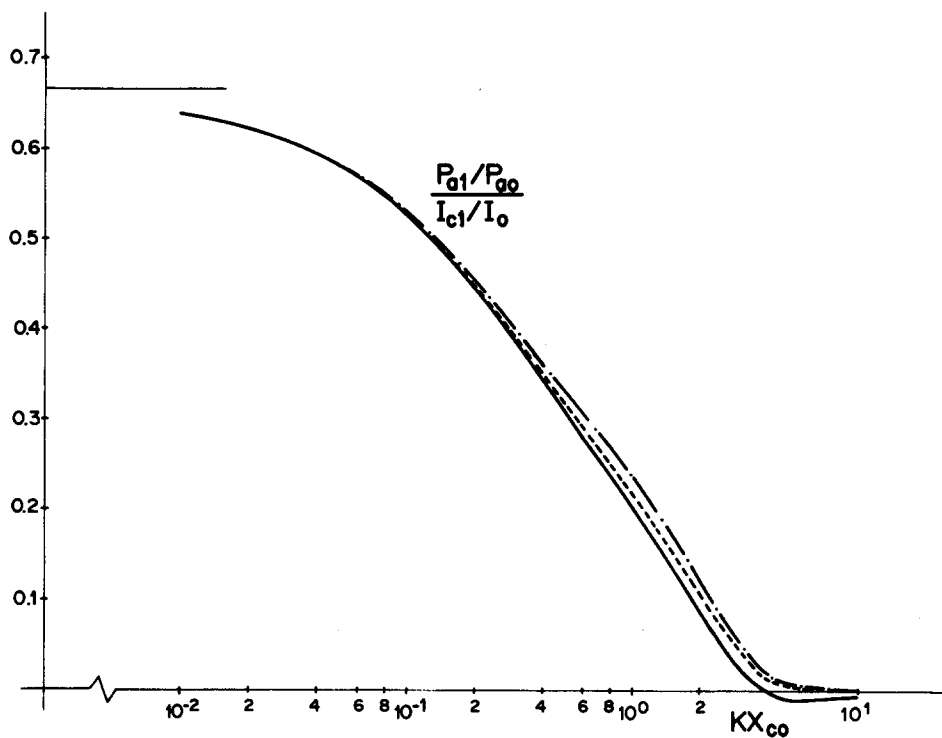


FIG. 2. Fractional perturbed ablation pressure normalized with fractional perturbed absorbed intensity; — $\Lambda_c = 0$, --- $\Lambda_c = 10$, — $\Lambda_c = \infty$ [Λ_c defined in Eqs. (15a) and (15b)]. Unperturbed distance between ablation and critical surfaces = x_{c0} , wavenumber = k .

arises from the sonic character of the zeroth-order solution at x_{c0} ; one of the modes of the linear system for the first-order fluid variables is singular there.

(7) and (8) Variables bounded as $s/x_{c0} \rightarrow \infty$; both system (13) and Eq. (14) have one homogeneous unbounded mode.

To carry out the integration, the behavior of the modes near $s = 0$ (where $n \rightarrow \infty$) and on both sides of x_{c0} (where

density and velocity gradients are singular) had to be determined. Since explicit formulas for n_0 , T_0 , and dv_0/ds in terms of v_0 are available from Eqs. (9)–(11), we used v_0 instead of s as an independent variable. For the underdense region, straightforward integration off the critical surface was sufficient to accurately determine x_{c1} and other overall quantities, although fluid profiles could only be roughly determined.

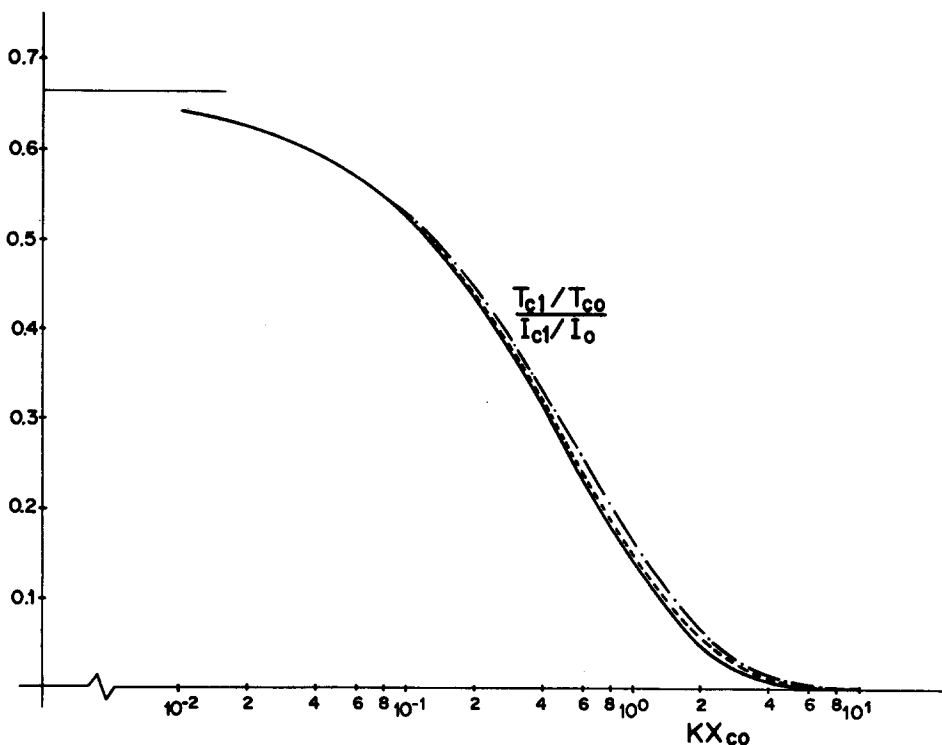


FIG. 3. Normalized perturbed critical temperature; other symbols as in Fig. 2.

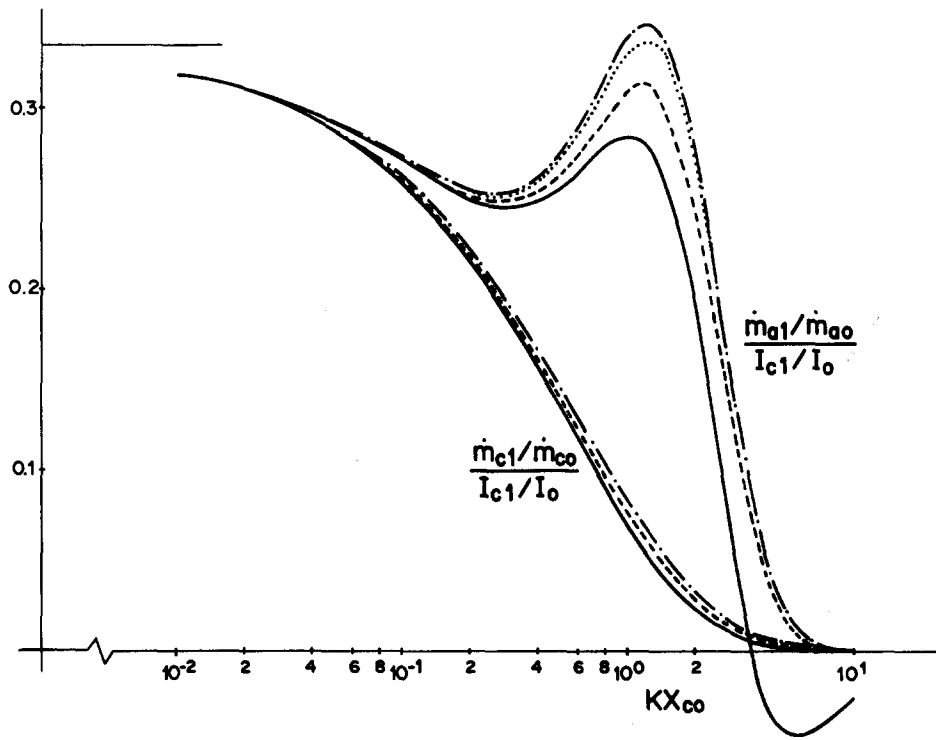


FIG. 4. Normalized perturbed mass flow at ablation (\dot{m}_a) and critical (\dot{m}_c) surfaces; $\cdots \Lambda_c = 0$, — results neglecting entirely the right-hand side of Eq. (13d), other symbols as in Fig. 2.

Two dimensionless parameters determine the solution: kx_{c0} and Λ_c . This second number characterizes current (low Λ_c) and magnetic (high Λ_c) effects, and is essentially the ratio of the electron mean free path to laser wavelength, both assumed small compared with x_{c0} . In fact, the results of Sec. III show that the mean free path is indeed less than x_{c0} by a factor of order $(m_e/\bar{m})^{1/2}$.

V. DISCUSSION

Figures 2–5 show the results (conveniently normalized) for perturbed ablation pressure P_{a1} , critical temperature T_{c1} (and critical pressure), mass flow rates \dot{m}_{a1} , \dot{m}_{c1} at ablation and critical surfaces, and location x_{c1} of critical density, respectively; e.g., Fig. 2 shows $(P_{a1}/P_{a0})/(I_{c1}/$

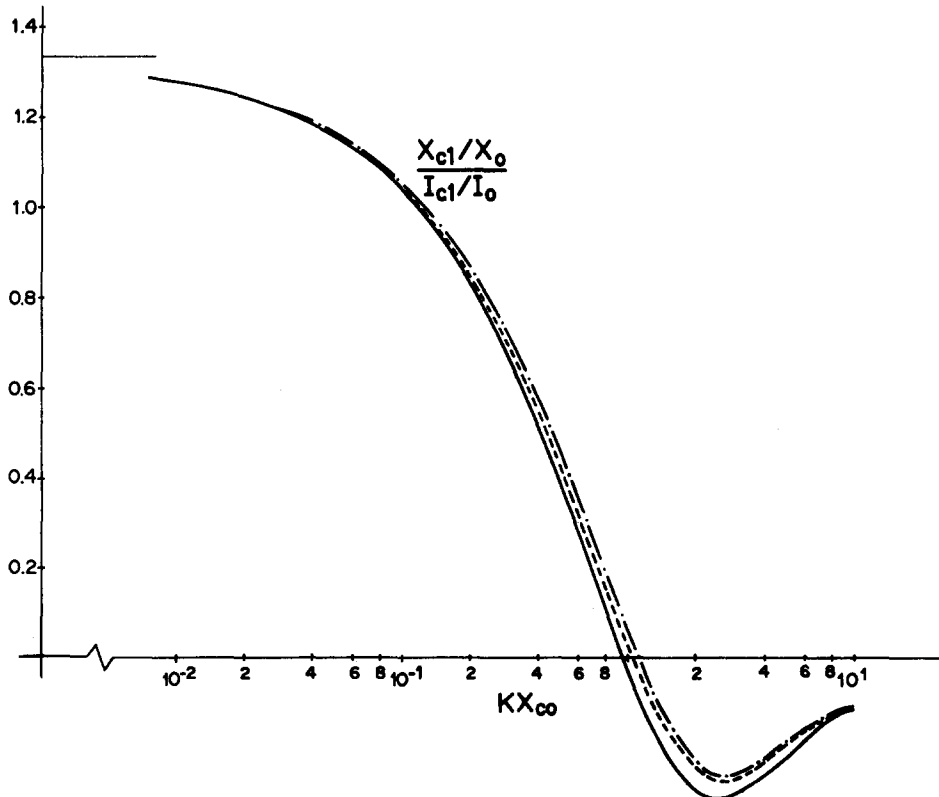


FIG. 5. Normalized perturbed distance between ablation and critical surfaces; other symbols as in Fig. 2.

I_{c0}). Values of $\Lambda_c = 0, 10, \infty$ for kx_{c0} arbitrary were considered. Because of the coordinate strain introduced in Sec. IV, T_{c1} in Fig. 3 represents the difference between the exact T at the exact x_c and T_{c0} . (Here, of course, "exact" means "to zero and first order.") We note the following features.

(i) The amount of smoothing increases as Λ_c becomes larger, but this dependence is weak. There is no practical difference between results for $\Lambda_c = 0$ and those obtained by simply dropping the entire right-hand side of (13d) (except in \dot{m}_{a1} , for which we show both). Low Λ_c values (and dominant current effects) correspond to short wavelength λ because the mean free path varies as $\lambda^{14/3}$; current effects are also favored in the overdense part of the layer [see Eqs. (15)].

(ii) As expected from the laws in Sec. III ($P_{c0} \sim T_{c0} \sim \dot{m}_{c0}^2 \sim x_{c0}^{1/2} \sim I_{c0}^{2/3}$), the ordinates in Figs. 2–5 go to $\frac{2}{3}$, $\frac{1}{3}$, $\frac{1}{3}$, and $\frac{1}{3}$, respectively, as $kx_{c0} \rightarrow 0$. Smoothing, which is practically nil at $kx_{c0} \sim 10^{-2}$, is important at $kx_{c0} \sim 10^{-1}$.

(iii) The mass ablation rate \dot{m}_{a1} exhibits a substantial peak at $kx_{c0} \sim 1$ (Fig. 4), and shows no noticeable smoothing. The effect is due to the overdense lateral mass flow (the difference between \dot{m}_{a1} and \dot{m}_{c1}) which must obviously vanish for both $kx_{c0} \rightarrow 0$ and $kx_{c0} \rightarrow \infty$, and peak somewhere. The mass ablation rate must sustain both that mass "loss" and the mass flow at $s = x_{c0}$.

(iv) The ratio x_{c1}/I_{c1} is negative beyond kx_{c0} around unity (Fig. 5) (note similar, weaker effects on \dot{m}_{a1} and P_{a1}). At such wavenumbers there is a substantial energy loss caused by the lateral flux. This requires a larger overdense temperature gradient normal to the target. If the increased critical temperature is unable to sustain that gradient, the critical surface moves closer to the target.

The analysis by Bell and Epperlein⁷ is similar to ours: weak nonuniformity, absorption at n_c , and conduction within a quasisteady layer (deflagration regime). The results, however, were quite different. As previously mentioned, the term $-T\mathbf{u} \cdot \nabla n$ was missing from their equations, but this

cannot account for the differences because the current and magnetic effects were weak [their results correspond to either dropping the right-hand side of (13d) or taking $\Lambda_c \rightarrow \infty$]. The root of the differences lies in the character of the zeroth-order solution they considered. They set the isothermal Mach number at the critical surface, M_{c0} , arbitrarily equal to 0.95, with fluid variables remaining constant throughout the underdense part of the layer. As shown in Sec. III, however, if absorption only occurs at n_c , it is impossible, in general, to match such a layer to the three-dimensional and/or unsteady flow that must exist outside (this flow was not discussed in Ref. 7).

Yet, if absorption by inverse bremsstrahlung is substantial, the outside flow is not isentropic, and the argument of Sec. III does not apply. This situation was recently studied for uniform irradiation.¹⁴ Two dimensionless parameters appeared in the study: the fractional absorption α at the critical surface, and a number characterizing inverse bremsstrahlung (\hat{V} for dominant spherical effects, \hat{U} for dominant unsteady effects). For small values of, say, \hat{V} we have the case considered in this paper: $M_{c0} = 1$ and $M_{\infty 0} = (\frac{2}{3})^{1/2}$, $M_{\infty 0}$ being the isothermal Mach number of the exit of the layer (the exit *isentropic* Mach number is then equal to unity). As \hat{V} is increased, M_{c0} remains unity while $M_{\infty 0}$ starts to decrease until it reaches M_{c0} . For larger \hat{V} , M_{c0} and $M_{\infty 0}$ decrease together, and at some stage the situation of Ref. 7 is attained.¹⁴ (For $M_{c0} = M_{\infty 0} < 1$, condition 6 of Sec. IV disappears, but then *two* fluid modes are unbounded at infinity so that the total number of boundary conditions does not change.) Note that in this case nonthermal smoothing outside the layer, caused by nonuniform bremsstrahlung absorption, must be considered, since the analysis of just the layer is *not* sufficient, as it is in our problem.

As an example of the differences mentioned above, Fig. 6 shows x_{c1} for the problem of Bell and Epperlein; that quantity was not given in Ref. 7. Note that x_{c1}/I_{c1} is negative for all wavenumbers, and that $(x_{c1}/x_{c0})/(I_{c1}/I_{c0})$ does not ap-

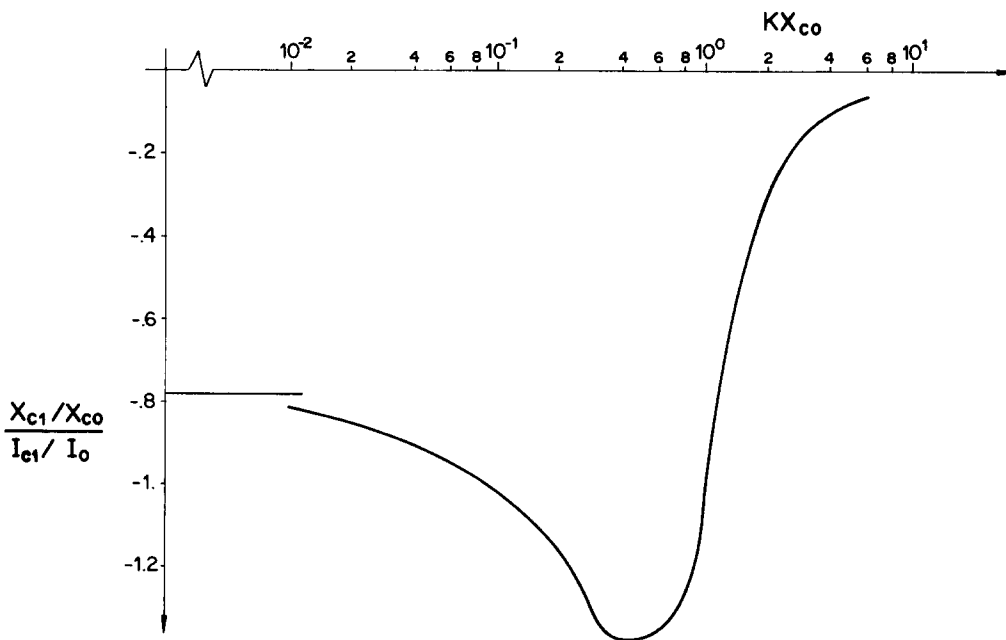


FIG. 6. Normalized perturbed distance between ablation and critical surfaces for the problem of Ref. 7; results neglecting the right-hand side of Eq. (13d).

proach $\frac{1}{3}$ when $kx_{c0} \rightarrow 0$, as in Fig. 5. The ordinate at the origin in both Figs. 5 and 6 [and similarly for $(P_{a1}/P_{a0})(I_{c1}/I_{c0})$, etc.] may be obtained analytically by systematically using the limit $kx_{c0} \rightarrow 0$ in the analysis of the solution. For instance, for $M_{c0} = 1$, $1 < M_{\infty 0} < (\frac{1}{3})^{1/2}$ we obtain

$$\frac{x_{c1}/x_{c0}}{I_{c1}/I_{c0}} = \frac{4(M_{\infty 0}^2 - 1)(M_{\infty 0}^2 + 5)}{3M_{\infty 0}^4 + 6M_{\infty 0}^2 - 5},$$

yielding $\frac{1}{3}$ at $M_{\infty 0}^2 = \frac{1}{3}$ (Fig. 5). For $M_{c0} = M_{\infty 0} < 1$, a somewhat more involved expression is obtained; it is easily verified that for $M_{c0} = M_{\infty 0} \simeq 0.95$, $(x_{c1}/x_{c0})/(I_{c1}/I_{c0}) \simeq -0.77$, as in Fig. 6.

Results from the present work are clearly only indicative; in particular it is unrealistic to neglect inverse bremsstrahlung if conduction is restricted to a thin layer in the corona (long, low-intensity, short-wavelength pulses). Extensions from our work to other regimes might include, besides inverse bremsstrahlung, spherical and unsteady effects mixed up with the smoothing, nonclassical transport, light refraction (these results will be published in a future paper), and possible instabilities tied up with smoothing processes. The richness of coronal regimes (and corresponding zero order, symmetric flows) should lead to a variety of results that are of interest for the direct-drive approach.

ACKNOWLEDGMENTS

This work was supported by the Comisión Asesora de Investigación Científico-Técnica of Spain (Project 1983/82) and the Instituto de Fusión Nuclear.

- ¹J. H. Gardner and S. E. Bodner, *Phys. Fluids* **29**, 2672 (1986); J. Grun, M. E. Emery, C. K. Mauka, T. N. Lee, E. A. McLean, A. Mostovich, J. Stamper, S. Bodner, and S. P. Obenschain, *Phys. Rev. Lett.* **58**, 2672 (1987).
- ²J. Nuckolls, L. Wood, A. Thiessen, and G. Zimmerman, *Nature* **239**, 139 (1972).
- ³K. A. Brueckner and S. Jorna, *Rev. Mod. Phys.* **46**, 325 (1974).
- ⁴J. H. Gardner and S. E. Bodner, *Phys. Rev. Lett.* **47**, 1137 (1981).
- ⁵W. M. Manheimer, D. G. Colombant, and J. H. Gardner, *Phys. Fluids* **25**, 1644 (1982).
- ⁶R. G. Evans, *Laser Part. Beams* **3**, 273 (1985).
- ⁷A. R. Bell and E. M. Epperlein, *Plasma Phys. Controlled Fusion* **28**, 897 (1986).
- ⁸J. R. Sanmartín, J. Sanz, and J. A. Nicolás, *Phys. Lett. A* **124**, 81 (1987).
- ⁹S. I. Braginskii, in *Reviews of Plasma Physics* (Consultants Bureau, New York, 1965), Vol. 1, p. 205.
- ¹⁰J. R. Sanmartín, J. L. Montañés, J. Sanz, and R. Ramis, *Plasma Phys. Controlled Fusion* **27**, 983 (1985).
- ¹¹J. R. Sanmartín, J. L. Montañés, and A. Barrero, *Phys. Fluids* **26**, 2754 (1983).
- ¹²J. Sanz, A. Liñán, M. Rodríguez, and J. R. Sanmartín, *Phys. Fluids* **24**, 2098 (1981).
- ¹³J. R. Sanmartín and A. Barrero, *Phys. Fluids* **21**, 1957 (1978); A. Barrero and J. R. Sanmartín, *Plasma Phys.* **22**, 617 (1980).
- ¹⁴J. A. Nicolás and J. R. Sanmartín, *Plasma Phys. Controlled Fusion* **27**, 279 (1985); J. A. Nicolás, *ibid.* **28**, 1441 (1986).